

On SUGRA description of boost-invariant conformal plasma at strong coupling ¹

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Abstract. We study string theory duals of the expanding boost invariant conformal gauge theory plasmas at strong coupling. The dual supergravity background is constructed as an asymptotic late-time expansion, corresponding to equilibration of the gauge theory plasma. The absence of curvature singularities in the first few orders of the late-time expansion of the dual gravitational background unambiguously determines the equilibrium equation of the state, and the shear viscosity of the gauge theory plasma. While the absence of the leading pole singularities in the gravitational curvature invariants at the third order in late-time expansion determines the relaxation time of the plasma, the subleading logarithmic singularity can not be canceled within a supergravity approximation. We comment on the possible interpretations of this singularity.

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The best understood example of gauge theory/string theory correspondence of Maldacena [2] is an equivalence between $\mathcal{N} = 4$ $SU(N)$ supersymmetric Yang-Mills (SYM) theory and type IIB string theory on $AdS_5 \times S^5$. In the planar ('t Hooft) limit and at weak effective coupling $\lambda \equiv g_{YM}^2 N \ll 1$, perturbative gauge theory description is valid. When the 't Hooft coupling is large, $\lambda \gg 1$, the gauge theory is best described by type IIB supergravity on $AdS_5 \times S^5$.

The main motivation of this work is to use string theory (in a context of Maldacena correspondence) as a guiding principle in constructing *Non-equilibrium Quantum Field Theory*. As a first step in this direction we would like to use string theory to formulate dissipative relativistic theory of conformal fluids.

In the next section we describe boost-invariant expansion of conformal fluids from the phenomenological perspective. We start with the ideal CFT fluid dynamics, then discuss the first-order dissipative CFT fluid dynamics. We explain the reason of going beyond the first-order hydrodynamics and outline the Müller-Israel-Stewart transient theory [3, 4]. In section II we review aspects of gauge theory/string theory correspondence relevant for the discussion of a CFT plasma expansion. Specifically, we emphasize nonsingularity of the background geometry as a guiding principle to determine correct physics. In section III we review Janik-Peschanski proposal² [5] for string theory dual to

¹ Based on [1].

² See also [6, 7, 8].

boost-invariant plasma expansion. We discuss successes of the proposal (correct equation of state, shear viscosity, relaxation time), and explain singularities in the supergravity approximation. We conclude in section IV with the possible interpretation of the singularities and outline future directions.

I: BOOST-INVARIANT EXPANSION OF THE CONFORMAL FLUIDS

Consider expansion of a CFT fluid (a gauge theory plasma) in the boost-invariant frame³. We convert Minkowski frame

$$ds_4^2 = -dx_0^2 + dx_\perp^2 + dx_3^2$$

into a frame with the boost-invariance along x_3 direction

$$x_0 = \tau \cosh y, \quad x_3 = \tau \sinh y$$

$$ds_4^2 = -d\tau^2 + \tau^2 dy^2 + dx_\perp^2.$$

We further assume

$$\varepsilon = \varepsilon(\tau), \quad P = P(\tau),$$

for a local energy density ε and a pressure P in a fluid.

The stress-energy tensor of an ideal fluid takes form

$$T_{\mu\nu} \equiv T_{\mu\nu}^{equilibrium} = (\varepsilon + P) u_\mu u_\nu + p \eta_{\mu\nu},$$

where u^μ is a local 4-velocity of the fluid, $u^2 = -1$. Scale invariance of a CFT fluid implies

$$T_\mu^\mu = 0 \quad \Rightarrow \quad \varepsilon = 3p,$$

so that the stress-energy conservation in the boost-invariant frame takes a simple form

$$\nabla_\mu T^{\mu\nu} = 0 \quad \Rightarrow \quad \frac{d\varepsilon}{d\tau} = -\frac{4}{3} \frac{\varepsilon}{\tau}.$$

Above equation describes evolution of the energy density of the expanding fluid at equilibrium. We find

$$\varepsilon \propto \tau^{-4/3}.$$

Since energy density is the only scale in a conformal fluid, dimensional arguments readily establish proper time scaling for the fluid temperature T , entropy density s , shear viscosity η , and relaxation time τ_π :

$$T \propto \varepsilon^{1/4} \propto \tau^{-1/3}, \quad \eta \propto s \propto T^3 \propto \tau^{-1}, \quad \tau_\pi \propto T^{-1} \propto \tau^{1/3}.$$

³ This is expected to be a correct description of a central region of QGP produced in ultra-relativistic collisions of heavy nuclei [9].

In the first-order dissipative CFT fluid dynamics the stress-energy tensor is phenomenologically modified from its equilibrium value by a symmetric, transverse and traceless tensor $\tau_{\mu\nu}$ constructed from the first derivatives of the 4-velocity u^μ :

$$T_{\mu\nu} = T_{\mu\nu}^{equilibrium} + \tau_{\mu\nu}, \quad \tau^{\mu\nu} = -\eta \left(\Delta^{\mu\alpha} \Delta^{\nu\beta} (\nabla^\mu u^\nu + \nabla^\nu u^\mu) - \frac{2}{3} \Delta^{\mu\nu} \Delta^{\alpha\beta} \nabla_\alpha u_\beta \right),$$

where η is the shear viscosity of the fluid⁴, and

$$\Delta^{\mu\nu} \equiv \eta^{\mu\nu} + u^\mu u^\nu.$$

In this case, the stress-energy conservation in the boost-invariant frame leads to

$$\frac{d\varepsilon}{d\tau} = -\frac{4}{3} \frac{\varepsilon}{\tau} + \frac{4\eta}{3\tau^2}.$$

From the equilibrium scaling for the boost-invariant expansion we see that the viscous correction becomes subdominant as $\tau \rightarrow \infty$. Indeed,

$$\left. \frac{\varepsilon}{\tau} \right|_{equilibrium} \propto \frac{\tau^{-4/3}}{\tau} \propto \tau^{-7/3}, \quad \left. \frac{\eta}{\tau^2} \right|_{equilibrium} \propto \frac{\tau^{-1}}{\tau^2} \propto \tau^{-9/3}.$$

Thus we expect approach to equilibrium in the boost-invariant frame to correspond to the late-time dynamics.

Why do we need to consider dissipative hydrodynamics beyond the first-order? It turns out that the first-order hydrodynamics allows for an acausal signal propagation. Though we do not cover causality violation in the first-order hydrodynamics here⁵, such a violation would be clear once we discuss the second-order (causal) dissipative CFT fluid dynamics of Müller-Israel-Stewart (MIS) [3, 4]. Basic equations of MIS theory in the boost-invariant frame for a CFT fluid take form

$$0 = \frac{d\varepsilon}{d\tau} + \frac{\varepsilon + P}{\tau} - \frac{1}{\tau} \Phi,$$

$$0 = \frac{d\Phi}{d\tau} + \frac{\Phi}{\tau_\pi} + \frac{1}{2} \Phi \left(\frac{1}{\tau} + \frac{1}{\beta_2} T \frac{d}{d\tau} \left(\frac{\beta_2}{T} \right) \right) - \frac{2}{3} \frac{1}{\beta_2} \frac{1}{\tau},$$

where τ_π is the relaxation time, Φ is related to the dissipative part of the stress-energy, and

$$\beta_2 = \frac{\tau_\pi}{2\eta}.$$

From the equilibrium scaling, $\tau \rightarrow \infty$ limit corresponds effectively to $\tau_\pi \rightarrow 0$ and the second-order hydrodynamics of MIS reduces to the first-order hydrodynamics. Since in

⁴ Scale invariance forbids nonzero bulk viscosity for a conformal fluid.

⁵ See [10] for a detailed discussion of this.

this limit the relaxation is effectively instantaneous, it is not surprising that causality in the first-order hydrodynamics is violated.

We would like to apply the second-order hydrodynamics to $\mathcal{N} = 4$ $SU(N)$ SYM plasma in the boost-invariant frame. Again, since local temperature is the only scale in the problem, we expect

$$\varepsilon(\tau) = \frac{3}{8}\pi^2 N^2 T(\tau)^4, \quad P(\tau) = \frac{1}{3}\varepsilon(\tau), \quad \eta(\tau) = A s(\tau) = A \frac{1}{2}\pi^2 N^2 T(\tau)^3$$

$$\tau_\pi(\tau) = r \tau_\pi^{Boltzmann}(\tau) = r \frac{3\eta(\tau)}{2P(\tau)}$$

where A is the ratio of shear viscosity to entropy density and r is the relaxation time in units of Boltzmann relaxation time. From MIS equations as $\tau \rightarrow \infty$:

$$T(\tau) = \frac{\Lambda}{\tau^{1/3}} \left(1 + \sum_{k=1}^{\infty} \frac{t_k}{(\Lambda \tau^{2/3})^k} \right), \quad \Phi(\tau) = \frac{2}{3}\pi^2 N^2 A \frac{\Lambda^3}{\tau^2} \left(1 + \sum_{k=1}^{\infty} \frac{f_k}{(\Lambda \tau^{2/3})^k} \right),$$

where Λ is an arbitrary scale and

$$t_k = t_k(A, r), \quad f_k = f_k(A, r).$$

Notice that within the second-order MIS theory, $\mathcal{N} = 4$ hydrodynamics is uniquely specified by two dimensionless parameters: $\{A, r\}$. On the other hand, to completely specify a non-equilibrium state of a system one needs an *infinite* set of parameters. These additional parameters will show up as higher-order transport coefficients in constitutive relations between non-equilibrium part of the stress-energy tensor $\tau^{\mu\nu}$ and higher-order gradients of the 4-velocity u^μ . It is very difficult to specify such constitutive relations from the purely phenomenological approach, *i.e.*, along the lines of MIS theory. But it is precisely in this respect that the gauge/string theory correspondence of Maldacena can be useful⁶.

II: SOME ASPECTS OF GAUGE/THEORY STRING THEORY CORRESPONDENCE

Maldacena correspondence is a duality between a gauge theory and a full string theory. However, the correspondence is useful when it is computationally tractable. Typically, this implies a truncation of the full string theory to its low-energy supergravity approximation. One has to be careful though: a particular supergravity truncation is not always consistent. Inconsistencies of the truncation are often reflected in singularities of the resulting supergravity backgrounds. We can distinguish three main cases:

- (A) \Rightarrow in some cases singularities of the supergravity backgrounds are simply an indication that (further) Kaluza-Klein truncation of the supergravity is incorrect, and including a finite number of additional SUGRA modes (doing consistent KK truncation) one

⁶ For recent developments in this direction see [11, 12, 13, 14].

obtains a smooth geometry (for example: a black hole solution on the singular conifold with self-dual fluxes [15, 16]);

- (B) \Rightarrow in some cases singularities of the supergravity backgrounds are expected to be resolved by including an infinite set of string theory α' corrections — from the gauge theory perspective this would imply that an infinite set of gauge theory operators (of increasingly high dimension) would develop a vacuum expectation value at strong coupling;

- (C) \Rightarrow in some cases singularities of the supergravity backgrounds are not expected to be resolved within the full string theory, as this would falsify gauge/string correspondence — string theory would predict a gauge theory phase which can not be realized physically (for example: a singularity of the Klebanov-Tseytlin geometry [17] is not expected to be resolved in string theory preserving both the supersymmetry and the chiral symmetry).

A singularity in a given supergravity truncation would imply correspondingly that:

- (A) \Rightarrow a consistency of supergravity would determine additional operators on the gauge theory side that would develop a VEV at strong coupling (for a finite temperature Klebanov-Tseytlin background, SUGRA is smooth once a $U(1)$ fiber inside $T^{1,1}$ is warped [16, 18] \Leftrightarrow a dim-6 operators of the thermal gauge theory plasma develops a VEV [19, 20]);

- (B) \Rightarrow SUGRA truncation is not useful;

- (C) \Rightarrow a phase of the gauge theory with prescribed symmetries simply does not exist (a singular Klebanov-Tseytlin solution is replaced with a smooth Klebanov-Strassler solution [21] where the chiral symmetry is broken).

III: JANIK-PESCHANSKI PROPOSAL FOR STRING THEORY DUAL TO BOOST INVARIANT EXPANSION

Given the symmetries of the expanding boost-invariant $\mathcal{N} = 4$ SYM plasma, the most general parity invariant truncation of the dual type IIB SUGRA takes form

$$ds_{10}^2 = e^{-2\alpha(\tau,z)} \left\{ \frac{1}{z^2} \left[-e^{2a(\tau,z)} d\tau^2 + e^{2b(\tau,z)} \tau^2 dy^2 + e^{2c(\tau,z)} dx_{\perp}^2 \right] + \frac{dz^2}{z^2} \right\} \\ + e^{6/5\alpha(\tau,z)} (dS^5)^2$$

for the Einstein frame metric;

$$F_5 = \mathcal{F}_5 + \star \mathcal{F}_5, \quad \mathcal{F}_5 = -4Q \omega_{S^5}, \quad \phi = \phi(\tau, z),$$

for the 5-form (Q is a constant related to the rank of the gauge group) and the dilaton. We set

$$Q = 1 \quad \Leftrightarrow \quad R_{AdS_5} = 1.$$

Asymptotically as $z \rightarrow 0$

$$\{a, b, c, \alpha, \phi\} \rightarrow 0,$$

to insure that the boundary geometry is that of the boost-invariant frame; however, we assume that

$$a(\tau, z) \sim \mathcal{O}(z^4) \neq 0,$$

so that the local energy density in the expanding plasma is nonzero.

We proceed then as follows:

\Rightarrow we try to construct a nonsingular geometry everywhere in the bulk, subject to the above boundary conditions;

\Rightarrow we evaluate the stress-energy tensor one-point correlation function

$$\langle T_{\mu\nu}(\tau) \rangle = \frac{N_c^2}{2\pi} \lim_{z \rightarrow 0} \frac{g_{\mu\nu}^{(5)}(\tau) - \eta_{\mu\nu}}{z^4};$$

\Rightarrow we extract from $\langle T_{\mu\nu}(\tau) \rangle$

$$\varepsilon(\tau), \quad p(\tau),$$

and interpret results in the framework of the dissipative relativistic fluid dynamics.

We saw before that near-equilibrium hydrodynamics corresponds to the late-time asymptotic expansion of the boost-invariant conformal plasma. Thus, it appears natural to assume [5, 7, 8]

$$a(\tau, z) = a\left(\tau, v \equiv \frac{z}{\tau^s}\right),$$

as well as for the remaining SUGRA modes; then, study background geometry as asymptotic expansion in τ while keeping the scaling variable v finite.

To leading order as $\tau \rightarrow \infty$, the absence of singularities in

$$\mathcal{J}^{[2]} \equiv \mathcal{R}_{\mu\nu\rho\lambda} \mathcal{R}^{\mu\nu\rho\lambda}, \quad v^4 \rightarrow 3_-,$$

requires

$$s = \frac{1}{3}.$$

Given the value of s , the asymptotic expansion for the warp factors of the 5-dim geometry takes form

$$a(\tau, v) = a_0(v) + \frac{1}{\tau^{2/3}} a_1(v) + \frac{1}{\tau^{4/3}} a_2(v) + \frac{1}{\tau^2} a_3(v) + \mathcal{O}(\tau^{-8/3}),$$

$$b(\tau, v) = b_0(v) + \frac{1}{\tau^{2/3}} b_1(v) + \frac{1}{\tau^{4/3}} b_2(v) + \frac{1}{\tau^2} b_3(v) + \mathcal{O}(\tau^{-8/3}),$$

$$c(\tau, v) = c_0(v) + \frac{1}{\tau^{2/3}} c_1(v) + \frac{1}{\tau^{4/3}} c_2(v) + \frac{1}{\tau^2} c_3(v) + \mathcal{O}(\tau^{-8/3}),$$

which leads to the late-time expansion of the Riemann tensor squared

$$\mathcal{J}^{[2]} = \mathcal{J}_0^{[2]}(v) + \frac{1}{\tau^{2/3}} \mathcal{J}_1^{[2]}(v) + \frac{1}{\tau^{4/3}} \mathcal{J}_2^{[2]}(v) + \frac{1}{\tau^2} \mathcal{J}_3^{[2]}(v) + \mathcal{O}(\tau^{-8/3}).$$

Let's assume first

$$\alpha(\tau, v) \equiv 0, \quad \phi(\tau, v) \equiv 0,$$

which on the gauge theory side implies that neither $\langle \text{Tr} F^2(\tau) \rangle$ (dual to a dilaton) nor the dim-8 operator (dual to a SUGRA scalar α) develop a VEV. We emphasize that this is an assumption which might or might not be correct — we use nonsingularity condition of the dual string (supergravity) description to test it.

We find (up to the second subleading order)

$$\varepsilon(\tau) = \left(\frac{N^2}{2\pi^2} \right) \frac{1}{\tau^{4/3}} \left\{ 1 - \frac{2\eta_0}{\tau^{2/3}} + \left(\frac{10}{3}\eta_0^2 + \frac{C}{36} \right) \frac{1}{\tau^{4/3}} + \dots \right\}.$$

Matching the gauge theory expansion for the energy density with that of the dual gravitational description we find

$$\Lambda = \frac{\sqrt{2}}{3^{1/4}\pi}, \quad A = \frac{3^{3/4}}{2^{3/2}\pi} \eta_0, \quad r = -\frac{11}{18} - \frac{1}{108} \frac{C}{\eta_0^2}.$$

As we alluded to before, further expansion on the SUGRA side will *define* higher order dissipative relativistic dynamics.

For generic values of $\{\eta_0, C\}$:

$$\left\{ \mathcal{J}_2^{[2]}(v), \mathcal{J}_3^{[2]}(v) \right\} = \mathcal{O} \left(\frac{1}{(3-v^4)^4} \right), \quad v^4 \rightarrow 3_-.$$

Tuning

$$\eta_0 = \frac{1}{2^{1/2}3^{3/4}}, \quad C = 2\sqrt{3} \ln 2 - \frac{17}{\sqrt{3}},$$

all pole singularities in $\{\mathcal{J}_2^{[2]}(v), \mathcal{J}_3^{[2]}(v)\}$ are removed. Mapping above values on the gauge theory side we find the ratio of the shear viscosity to the entropy density and the plasma relaxation time (within MIS phenomenological theory)

$$A = \frac{1}{4\pi}, \quad r = \frac{1}{3}(1 - \ln 2).$$

Notice that the shear viscosity agrees⁷ with earlier computations using equilibrium correlation functions [22].

However:

$$\mathcal{J}_3^{[2]} = \text{finite} + \left(8 \cdot 2^{1/2} \cdot 3^{3/4} \right) \ln(3-v^4), \quad v \rightarrow 3_-^{1/4}.$$

Thus it appears inconsistent to set α and/or the dilaton to zero; in fact weak coupling analysis suggests that there are instabilities in expanding plasma generating VEVs of various operators [27], in particular $\langle \text{Tr} F^2(\tau) \rangle$. A careful analysis shows that without

⁷ Such an agreement is lost once finite 't Hooft coupling corrections are taken into account [23, 24, 25]. The latter suggests that the commonly used $\mathcal{O}(\alpha'^3)$ string theory effective action is incomplete (see also [26]).

introducing additional pole curvature singularities one can turn on only the α mode to relevant order

$$\alpha(\tau, v) = \frac{1}{\tau^2} \alpha_3(v) + \mathcal{O}(\tau^{-8/3}),$$

$$\alpha_3 = \alpha_{3,0} \left(\left(\frac{1}{96v^4} + \frac{v^4}{864} \right) \ln \frac{3+v^4}{3-v^4} - \frac{1}{144} \right),$$

where $\alpha_{3,0}$ is a normalizable mode related to the VEV of the dim-8 operator in $\mathcal{N} = 4$ SYM plasma.

We find:

$$\mathcal{I}_3^{[2]} = \text{finite} + \left(8 \cdot 2^{1/2} \cdot 3^{3/4} + \frac{14}{3} \alpha_{3,0} \right) \ln(3-v^4), \quad v \rightarrow 3_-^{1/4},$$

but

$$\mathcal{R}_{\mu\nu} \mathcal{R}^{\mu\nu} = \text{finite} + \frac{1}{\tau^2} \frac{40}{3} \alpha_{3,0} \ln(3-v^4), \quad v \rightarrow 3_-^{1/4}.$$

So, logarithmic singularity can not be canceled within the SUGRA approximation.

We considered another model of a CFT plasma (specifically, Klebanov-Witten plasma [28]) which has an additional SUGRA mode and showed that logarithmic singularities both in

$$\mathcal{R}_{\mu\nu} \mathcal{R}^{\mu\nu} \quad \text{and} \quad \mathcal{R}_{\mu\nu\rho\lambda} \mathcal{R}^{\mu\nu\rho\lambda}$$

at the third subleading order can be canceled (Ricci scalar is nonsingular). However, new logarithmic singularities appear at the third order in higher curvature invariants such as

$$\mathcal{R}_{\mu_1\nu_1\lambda_1\rho_1} \mathcal{R}^{\mu_1\nu_1\lambda_2\rho_2} \mathcal{R}_{\mu_2\nu_2}{}^{\lambda_1\rho_1} \mathcal{R}^{\mu_2\nu_2}{}_{\lambda_2\rho_2},$$

as well as logarithmic singularities with different coefficients in

$$(\mathcal{R} \dots)^8, \quad (\mathcal{R} \dots)^{16},$$

and so on. Thus, we conclude that one needs an infinite set of fields to cancel singularities in gravitational description, corresponding to an infinite set of gauge invariant operators developing a VEV during boost-invariant expansion.

IV: INTERPRETATION OF SINGULARITIES AND FUTURE DIRECTIONS

We attempted to construct a string theory dual to strongly coupled conformal expanding plasmas in the Bjorken regime. In order to have computational control we truncated the full string theory to supergravity approximation, and focused on the well-established examples of the gauge/string dualities: we considered $\mathcal{N} = 4$ SYM and superconformal Klebanov-Witten gauge theories. We used nonsingularity of the dual gravitational backgrounds as a guiding principle to identify gauge theory operators that would develop a vacuum expectation value during boost-invariant expansion of the plasma. Truncation

to a supergravity sector of the string theory (along with parity invariance in the Bjorken frame) severely restricts a set of such operators. In the case of the $\mathcal{N} = 4$ SYM, there are only two such gauge invariant operators, while for the Klebanov-Witten plasma one has an additional operator. We constructed supergravity dual as a late-time asymptotic expansion and demonstrated that the gravitational boundary stress energy tensor expectation value has exactly the same asymptotic late-time expansion as predicted by Müller-Israel-Stewart theory of transient relativistic kinetic theory for the boost-invariant expansion. As an impressive success of this approach, one recovers by requiring nonsingularity of the background geometry at leading and the first three subleading orders⁸ the equation of state for the plasma and its shear viscosity, in agreement with values extracted from the equilibrium correlation functions. Unfortunately, we showed that logarithmic singularity in the background geometry can not be canceled within the supergravity approximation. Moreover, given that the singularities appear to persist in arbitrary high order metric curvature invariants, we suspect that relaxing the constraint of parity invariance in the Bjorken regime would not help. Indeed, relaxing parity invariance would allow for only finite number of additional (massive) supergravity modes, which, as an example of Klebanov-Witten plasma demonstrates, would only allow to cancel logarithmic singularities in a finite number of additional metric curvature invariants.

Coming back to our classification of singularities in SUGRA truncations of the full string theory we see that for the string theory duals of the boost-invariant CFT plasmas expansion

- (A) \Rightarrow is not realized;
- (B) \Rightarrow can be realized; in this case, though SUGRA truncation is inconsistent, maybe the requirement of the cancellation of the pole singularities at low orders is a correct prescription to extract the second order transport coefficients (which are of relevance to RHIC); tantalizingly, we see hints of the universality of the relaxation time — further study of non-conformal models is needed;
- (C) \Rightarrow can be realized; in this case a SUGRA singularity might be an indication of a genuine singularity in the full string theory description — here, it is important to search for onset of instabilities/turbulence in the expanding plasma.

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⁸ At the third subleading order, logarithmic singularities of the metric curvature invariants remain.

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